

Pion form factor at spacelike momentum transfers from local-duality QCD sum rule

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(Dated: February 6, 2008)

We study the pion form factor in a broad range of spacelike momentum transfers within the local-duality version of QCD sum rules. We make use of the recently calculated two-loop double spectral density of the $\langle AVA \rangle$ correlator including $O(1)$ and $O(\alpha_s)$ terms, which allows us to give predictions for the pion form factor and to study the interplay between the nonperturbative and perturbative contributions to the pion form factor without any reference to the pion distribution amplitude. Our results demonstrate the dominance of the nonperturbative contribution to the form factor up to relatively large values of the momentum transfer: namely, the nonperturbative $O(1)$ term, which provides the $1/Q^4$ power correction, gives more than half of the pion form factor in the region $Q^2 \leq 20 \text{ GeV}^2$.

PACS numbers: 11.55.Hx, 12.38.Lg, 13.40.Gp

1. INTRODUCTION

The study of the interplay between perturbative and nonperturbative physics in exclusive processes, and, in particular, in the pion form factor which we discuss in this paper, has a long history. At asymptotically large Q^2 , the pQCD factorization formula [1] gives the pion form factor in terms of the scale-dependent pion distribution amplitude (DA) of leading twist $\phi_\pi(u, Q^2)$:

$$F_\pi(Q^2) = \frac{8\pi\alpha_s(Q^2)f_\pi^2}{9Q^2} \left| \int_0^1 du \frac{\phi_\pi(u, Q^2)}{u} \right|^2. \quad (1.1)$$

The DA, as obtained from the pQCD evolution equation, has the form

$$\phi_\pi(u, Q^2 \rightarrow \infty) = 6u(1-u). \quad (1.2)$$

Respectively, at asymptotically large Q^2 a direct prediction of pQCD reads [1, 2]

$$Q^2 F_\pi(Q^2) = 8\pi\alpha_s(Q^2)f_\pi^2. \quad (1.3)$$

Subleading logarithmic and power corrections to this formula should be taken into account at large but finite Q^2 . This is, however, a very difficult task. There are two competitive scenarios for the pion form factor at intermediate momentum transfers:

The first scenario (A) is based on the assumption that power corrections are negligible in the region $Q^2 \geq 3-5 \text{ GeV}^2$. The form factor is then given by the pQCD factorization formula (1.1) with the pion distribution amplitude at low normalization scale, which in this scenario turns out to have a double-humped “camel” shape with an enhanced end-point region [3], very different from its asymptotic form (1.2). This scenario, complemented by the analysis of Sudakov double logarithms [4], provides the basis for the perturbative QCD approach to form factors at intermediate momentum transfers.

In the second scenario (B), which we consider to be more realistic, the form factor is dominated by the nonperturbative contributions up to rather high values of Q^2 , with the perturbative contribution remaining relatively small [5, 6]. The end-point behaviour of the DA at a low normalization scale is then similar to that of the asymptotic DA (1.2). This scenario is supported by the fact that the soft contribution to the form factor alone can reproduce the pion form factor to a good accuracy for Q^2 up to several GeV^2 [6, 7, 8, 9, 10, 11].

In [12], making use of the constituent quark picture, and in [13], within light-cone sum rules, the pion form factor was analyzed by taking into account the nonperturbative $O(1)$ contribution and the radiative $O(\alpha_s)$ corrections. The form factor at intermediate Q^2 turned out to be sensitive to the details of the pion wave function — the Bethe-Salpeter wave function in [12] and the pion light-cone DA in [13]. Scenario B was favoured by these results.

Unfortunately, the data on the pion form factor for $Q^2 > 2 \text{ GeV}^2$ are not sufficiently precise, leaving room for speculations about the details of the pion DA at low normalization scale and, respectively, on the relative weights of the soft and the hard contributions to the pion form factor.

Therefore, it seems interesting to address the problem without a direct reference to the pion DA. The local-duality version of three-point QCD sum rules [6] provides this opportunity.

The local-duality sum rule is the Borel sum rule in the limit of an infinitely large Borel parameter. For the relevant choice of the pion interpolating current, the condensate contributions to the correlator vanish in this limit and the pion observables are given by dispersion integrals via the spectral densities of purely perturbative QCD diagrams. The integration region in the dispersion integrals is restricted to the pion “duality interval”.

In this Letter we apply a local-duality sum rule to the pion form factor, making use of the recently calculated two-loop double spectral density of the pion form factor for massless quarks [14, 15]. Such an approach has the following attractive features: (i) it is applicable in a broad range of momentum transfers starting from low to asymptotically large values, and (ii) it does not refer to the pion distribution amplitude. Therefore, it allows us to study in a relatively model-independent way the interplay between perturbative and nonperturbative dynamics in the pion form factor.

2. SUM RULE

We shall consider the pion form factor in the chiral limit of massless quarks and a massless pion. Let us recall well-known results for Borel sum rules: The sum rule for the pion decay constant is obtained from the OPE for the two-point function and reads [16]

$$f_\pi^2 = \frac{1}{\pi} \int_0^{s_0} ds \exp(-s/M^2) \rho(s) + \frac{\langle \alpha_s G^2 \rangle}{12\pi M^2} + \frac{176\pi\alpha_s \langle \bar{q}q \rangle^2}{81M^4} + \dots, \quad (2.1)$$

where $\rho(s) = \frac{1}{4\pi} (1 + \frac{\alpha_s}{\pi}) + O(\alpha_s^2)$ is the perturbative spectral density.

The Borel sum rule for the pion form factor is obtained from the OPE for the three-point function and reads [6, 7]

$$f_\pi^2 F_\pi(Q^2) = \Gamma(Q^2, M^2, M^2|s_0) + \frac{\langle \alpha_s G^2 \rangle}{24\pi M^2} + \frac{4\pi\alpha_s \langle \bar{q}q \rangle^2}{81M^4} \left(13 + \frac{Q^2}{M^2} \right). \quad (2.2)$$

Here, $\Gamma(Q^2, M^2, M^2|s_0)$ is the perturbative contribution, which is obtained by the following procedure: One calculates the double Borel transform of the $\langle AVA \rangle$ correlator

$$\Gamma(Q^2, M_1^2, M_2^2) = \frac{1}{\pi^2} \int ds_1 ds_2 \exp(-s_1/M_1^2) \exp(-s_2/M_2^2) \left[\Delta^{(0)}(Q^2, s_1, s_2) + \alpha_s \Delta^{(1)}(Q^2, s_1, s_2) \right], \quad (2.3)$$

and restricts the integration in the s_1 - s_2 plane to the pion duality region. One then sets $M_1^2 \rightarrow 2M^2$, $M_2^2 \rightarrow 2M^2$ and compares the two- and three-point sum rules for the same values of the Borel parameter M^2 .¹ The function $\Delta^{(0)}(Q^2, s_1, s_2)$ is well-known, whereas $\Delta^{(1)}(Q^2, s_1, s_2)$ was calculated only recently [15] for the case of massless quarks.² The explicit expressions can be found in [15].

The local-duality (LD) sum rules [6, 8, 19] correspond to the limit $M \rightarrow \infty$. A remarkable feature of this limit is the vanishing of the condensate contributions to the sum rules for the pion form factor and the decay constant. Assuming the duality region in the s_1 - s_2 plane to be a square of side s_0 , and denoting the duality interval in the sum rule for the decay constant by \bar{s}_0 , we obtain to α_s accuracy

$$f_\pi^2 F_\pi(Q^2) = \frac{1}{\pi^2} \int_0^{s_0} ds_1 \int_0^{s_0} ds_2 \left[\Delta^{(0)}(s_1, s_2, Q^2) + \alpha_s \Delta^{(1)}(s_1, s_2, Q^2) + O(\alpha_s^2) \right], \quad (2.4)$$

$$f_\pi^2 = \frac{1}{\pi} \int_0^{\bar{s}_0} ds \left[\rho^{(0)}(s) + \alpha_s \rho^{(1)}(s) + O(\alpha_s^2) \right] = \frac{\bar{s}_0}{4\pi^2} \left(1 + \frac{\alpha_s}{\pi} \right) + O(\alpha_s^2). \quad (2.5)$$

One should not be confused by the simplicity of these expressions: the complicated nonperturbative dynamics is now hidden in the effective continuum thresholds s_0 and \bar{s}_0 . Let us emphasize that the LD sum rules are predictive only if one knows, or fixes according to some criteria, the effective continuum thresholds.

Some comments on the dispersion representations for $f_\pi^2 F_\pi$ and f_π^2 are in order:

¹ Notice that such a procedure of comparing two- and three-point sum rules finds a natural physical explanation within the correspondence between sum rules and the constituent quark picture observed in [17].

² Another case in which the radiative corrections to the double spectral density of the three-point function have been calculated is the case of one massless and one infinitely heavy quark [18].

1. Whereas the single dispersion representation for the decay constant f_π^2 is well-defined, the double dispersion representation for $f_\pi^2 F_\pi$, even in the LD limit, has at least two essential ambiguities:

a. The shape of the duality region in the s_1 – s_2 plane: the simplest choice is a square, but any other region symmetric under $s_1 \leftrightarrow s_2$ may be also possible.

b. Nothing forbids the upper boundary of the duality region s_0 from being Q^2 -dependent, and additional assumptions to fix $s_0(Q^2)$ are necessary. Arguments in favour of choosing the parameters in two- and three-point sum rules constant and equal to each other were given in [8]. Let us see what happens if we choose the same constant value $\bar{s}_0 = s_0$ in the sum rules (2.4) and (2.5), and substitute the sum rule (2.5) instead of f_π^2 into (2.4):

$$F_\pi(Q^2) = \frac{\frac{1}{\pi^2} \int_0^{s_0} ds_1 \int_0^{s_0} ds_2 \left[\Delta^{(0)}(s_1, s_2, Q^2) + \alpha_s \Delta^{(1)}(s_1, s_2, Q^2) \right]}{\frac{1}{\pi} \int_0^{s_0} ds \left[\rho^{(0)}(s) + \alpha_s \rho^{(1)}(s) \right]}, \quad s_0 = \frac{4\pi^2 f_\pi^2}{1 + \alpha_s/\pi}. \quad (2.6)$$

The LD form factor given by (2.6) has the following interesting properties [20]:

(i) It satisfies the normalization condition $F_\pi(Q^2 = 0) = 1$ due to the vector Ward identity which relates the spectral density of the self-energy diagram and the double spectral density of the triangle diagram at zero momentum transfer:

$$\lim_{Q^2 \rightarrow 0} \Delta^{(i)}(s_1, s_2, Q^2) = \pi \rho^{(i)}(s_1) \delta(s_1 - s_2), \quad \rho^{(0)}(s) = \frac{1}{4\pi}, \quad \rho^{(1)}(s) = \frac{1}{4\pi^2}. \quad (2.7)$$

Clearly, for consistency one should then take into account the radiative corrections to the same order in two- and three-point correlators.

(ii) Making use of the explicit expressions for $\Delta^{(i)}$, one obtains

$$\lim_{Q^2 \rightarrow \infty} \Delta^{(0)}(s_1, s_2, Q^2) = \frac{3(s_1 + s_2)}{2Q^4}, \quad (2.8)$$

$$\lim_{Q^2 \rightarrow \infty} \Delta^{(1)}(s_1, s_2, Q^2) = \frac{1}{2\pi Q^2}. \quad (2.9)$$

Substituting these expressions into (2.6), one finds at large Q^2 :

$$F_\pi(Q^2) = \frac{8\pi f_\pi^2 \alpha_s}{Q^2} + \frac{96\pi^4 f_\pi^4}{Q^4} + O(\alpha_s f_\pi^4 / Q^4) + O(\alpha_s^2). \quad (2.10)$$

We find quite remarkable that the exact $\Delta^{(1)}$ leads to the correct pQCD (up to the running of α_s) large- Q^2 asymptotics of the pion form factor obtained from the LD sum rule (2.6). Let us explain this important point: Whereas, e.g., the normalization of the pion form factor (2.6) at $Q^2 = 0$ is the consequence of the Ward identity, we do not see any *rigorous* condition which would guarantee the correct large- Q^2 behaviour when using the same value of the pion duality intervals in two- and three-point correlators. We find this to be a strong argument in favour of the universality of the pion duality interval.

(iii) The $O(1)$ contribution, shown in Fig. 1a, was calculated in [6]:

$$I_0(Q^2, s_0) = \int_0^{s_0} ds_1 \int_0^{s_0} ds_2 \Delta^{(0)}(s_1, s_2, Q^2) = \frac{s_0}{4} \left(1 - \frac{1 + 6s_0/Q^2}{(1 + 4s_0/Q^2)^{3/2}} \right). \quad (2.11)$$

The explicit expression for the $O(\alpha_s)$ contribution

$$I_1(Q^2, s_0) = \int_0^{s_0} ds_1 \int_0^{s_0} ds_2 \Delta^{(1)}(s_1, s_2, Q^2) \quad (2.12)$$

was obtained only recently [15]. Before that, it was proposed to use instead of the unknown integral the simplest Ansatz [19]

$$I_1(Q^2, s_0) \rightarrow \frac{s_0}{4\pi} \frac{1}{1 + Q^2/2s_0}, \quad (2.13)$$

which reproduces the value of the integral at $Q^2 = 0$, fixed by the Ward identity, and its asymptotic behaviour according to (2.9). Fig. 1b compares the formula (2.13) with the result of the exact calculation: as one can see, the proposed formula underestimates the exact $I^{(1)}$ by more than 20% in the broad range of practical relevance $Q^2 = 1$ – 30 GeV².

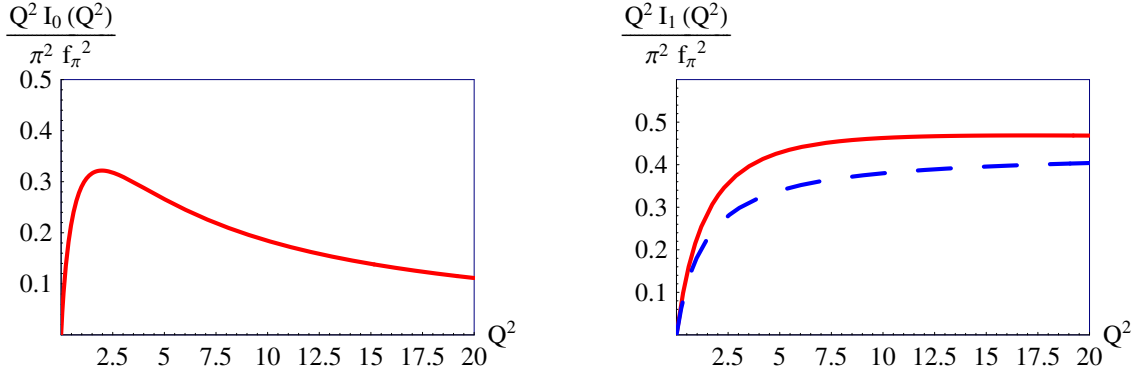


Fig. 1: $\frac{Q^2}{\pi^2 f_\pi^2} I_0(Q^2, s_0)$ (a) and $\frac{Q^2}{\pi^2 f_\pi^2} I_1(Q^2, s_0)$ (b) for $s_0 = 4\pi^2 f_\pi^2$. Solid (red) line: exact result, dashed (blue) line: Ansatz (2.13).

2. There are obvious problems with the application of this sum rule at small $Q^2 \leq 1 \text{ GeV}^2$:

First, the OPE for the three-point correlator was obtained in the region where all three external variables $|p_1^2|$, $|p_2^2|$, and Q^2 are large. Therefore, the sum rule cannot be directly applied at small Q^2 , although the expression (2.6) leads to the correct normalization of the form factor. Additional contributions appear at small Q^2 [21] which prevent from giving unambiguous predictions in this region. Of course, allowing for a Q^2 -dependent value $s_0(Q^2)$ in the sum rule (2.4), we can formally extend the formula also to lower Q^2 and apply it starting from $Q^2 = 0$, but as we have noticed above, in this case the sum rule loses its predictivity. A technical indication that the LD sum rule (2.6) cannot be applied at very small Q^2 is the presence of the terms $\sim \sqrt{Q^2}$ [see (2.11)] leading to an infinite value of the pion radius.

Second, the spectral density $\Delta(s_1, s_2, Q^2)$ contains the terms $O(1)$ and $O(\alpha_s)$, whereas higher powers are unknown. Since the coupling constant α_s is not small in the soft region, our spectral density is not sufficient for application to the form factor at $Q^2 \leq 1 \text{ GeV}^2$.

3. In order to apply the obtained formulas for large Q^2 , higher-order radiative corrections, leading to the running of α_s , should be taken into account. Such an accuracy is beyond our two-loop calculation; nevertheless, a self-consistent expression for the form factor applicable for all $Q^2 > 0$ may be written as

$$F_\pi(Q^2) = \frac{1}{f_\pi^2 \pi^2} \int_0^{s_0(Q^2)} ds_1 \int_0^{s_0(Q^2)} ds_2 \left[\Delta^{(0)}(s_1, s_2, Q^2) + \alpha_s(Q^{*2}) \Delta^{(1)}(s_1, s_2, Q^2) \right], \quad (2.14)$$

where the scale Q^{*2} in the argument of α_s is related to Q^2 (see the discussion in [22]) and $s_0(Q^2)$ satisfies the boundary conditions³

$$s_0(Q^2 = 0) = \frac{4\pi^2 f_\pi^2}{1 + \alpha_s(0)/\pi}, \quad s_0(Q^2 = \infty) = 4\pi^2 f_\pi^2. \quad (2.15)$$

If the effective threshold $s_0(Q^2)$ satisfies these relations, the form factor is normalized to $F_\pi(0) = 1$ and reproduces the pQCD asymptotic behaviour at $Q^2 \rightarrow \infty$.

In the following, we shall set the scale $Q^{*2} = Q^2$: in the region $Q^2 \geq 1 \text{ GeV}^2$, $\alpha_s(Q^2)$ is a slowly varying function of Q^2 (Fig. 2a); therefore, the precise setting of the scale makes very little difference. We shall use the following appealing parametrization of $s_0(Q^2)$, obviously satisfying (2.15):

$$s_0(Q^2) = \frac{4\pi^2 f_\pi^2}{1 + \alpha_s(Q^2)/\pi}. \quad (2.16)$$

Before turning to the numerical analysis, we would like to draw the reader's attention to the following observation: An essential feature of the form factor obtained from the three-point sum rule is the full cancellation of the double logarithmic terms. The proof of this general property of the color-neutral three-point Green functions in QCD can

³ In [19] it was argued that s_0 is the relevant scale of α_s in the LD sum rules for the decay constant and for the form factor at $Q^2 = 0$. For our discussion this subtlety is irrelevant so we somewhat symbolically use the notation $\alpha_s(0)$.

be found in [3] (see also [23]); the cancellation of double logs was checked in explicit two-loop calculations for various quark currents in [14, 15]. In contrast to this result, the pion form factor obtained from the light-cone sum rule contains double log terms [13]. This discrepancy requires a clarification. Presumably, higher-twist contributions, which were not taken into account in [13], but which are in general not suppressed compared to the lower-twist contributions [24] play a crucial role here.

3. NUMERICAL RESULTS

For numerical estimates, we make use of the three-loop running $\alpha_s(Q^2)$ (Fig. 2a). The corresponding $s_0(Q^2)$ is shown in Fig. 2b. Notice that it is a slowly varying function in the region $Q^2 \geq 1 \text{ GeV}^2$, where we apply the LD sum rule.

The pion form factor is shown in Fig. 3a. The $O(1)$ and $O(\alpha_s)$ terms, separately, are given in Fig. 3b. It should be noticed that the $O(1)$ term providing the $1/Q^4$ power correction at large Q^2 , dominates the form factor at low Q^2 , and still gives 50% at $Q^2 = 20 \text{ GeV}^2$. The $O(\alpha_s)$ term gives more than 80% of the form factor only above $Q^2 = 100 \text{ GeV}^2$. Such a pattern of the pion form factor behaviour has been conjectured many times in the literature; we now obtain this behaviour in an explicit calculation. The earlier analyses of the pion form factor in a broad range of momentum transfers [12, 13, 22, 26] are consistent with the results reported in Fig. 3 within about 20% accuracy.

Fig. 4 presents the ratio of the $O(1)$ and the $O(\alpha_s)$ contributions to the pion form factor vs Q^2 for different models of the effective continuum thresholds. One can clearly see that the ratio is mainly determined by the corresponding double spectral densities $\Delta^{(0)}$ and $\Delta^{(1)}$, whereas its sensitivity to the effective continuum threshold is rather weak.

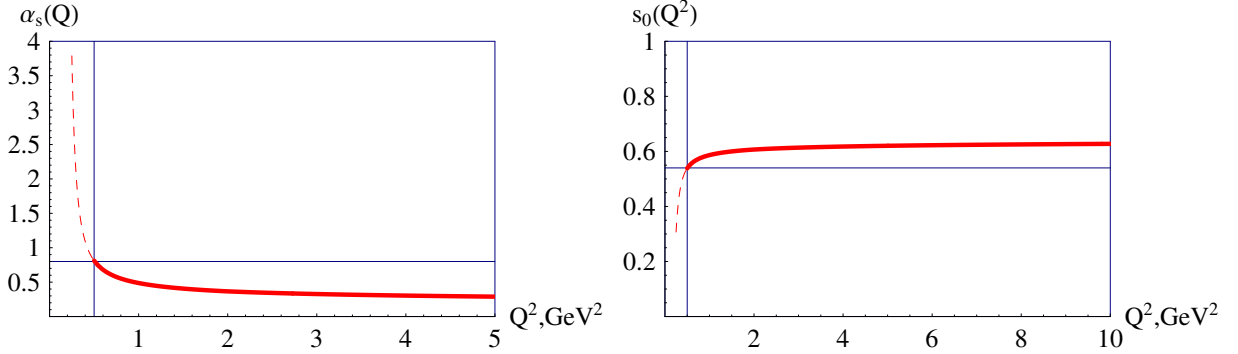


Fig. 2: The perturbative $\alpha_s(Q)$ (a) and the corresponding effective threshold $s_0(Q^2)$ (b) given by (2.15). Dashed lines show these quantities outside our working region.

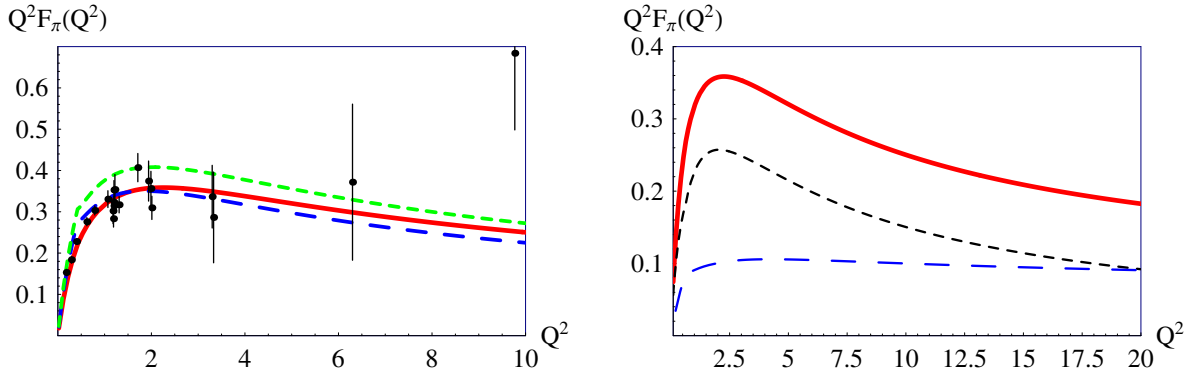


Fig. 3: The pion form factor at $Q^2 \geq 0.5 \text{ GeV}^2$. Experimental data from [25]. Solid (red) line: the result of the calculation according to (2.14). (a) Short-dashed (green) line: the form factor obtained with constant $s_0 = 0.65 \text{ GeV}^2$; long-dashed (blue) line: $s_0 = 0.6 \text{ GeV}^2$. (b) Short-dashed (black) line: the $O(1)$ contribution, long-dashed (blue) line: the $O(\alpha_s)$ contribution.

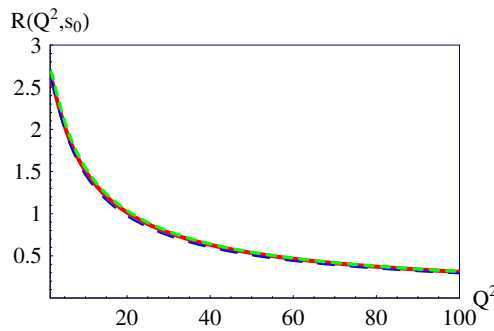


Fig. 4: The ratio of the $O(1)$ and $O(\alpha_s)$ contributions to the pion form factor $R(Q^2, s_0) = I_0(Q^2, s_0)/[\alpha_s(Q^2)I_1(Q^2, s_0)]$. Solid (red) line: the result of the calculation with the effective continuum threshold $s_0(Q^2)$ (2.16), short-dashed (green) line: $s_0 = 0.65 \text{ GeV}^2$, long-dashed (blue) line: $s_0 = 0.6 \text{ GeV}^2$.

4. DISCUSSION AND CONCLUSIONS

We have presented the analysis of the pion form factor in a broad range of spacelike momentum transfers making use of the local-duality sum rule. This is the first analysis which takes into account both the leading order $O(1)$ contribution to the pion form factor and the recently calculated first-order $O(\alpha_s)$ radiative correction. These ingredients are crucial for the possibility to consider the form factor in a broad range of Q^2 and to study the transition from the nonperturbative to the perturbative region.

Let us summarize the essential ingredients, the uncertainties, and the lessons to be learnt from our analysis:

- *The double spectral density of the spectral representation for the form factor:* We have good control over the spectral density — we have included the exact $O(1)$ and $O(\alpha_s)$ terms, and omit the (unknown) $O(\alpha_s^2)$ terms, which are expected to contribute less than 10% at $Q^2 > 1 \text{ GeV}^2$. [The inclusion of the $O(\alpha_s^2)$ terms in the spectral density would lead to a corresponding modification of the effective continuum threshold, with the net effect upon the form factor of only a few percent.]
- *The model for the effective continuum threshold:* This very quantity determines to a great extent the value of the form factor extracted from the sum rule. The possibility to fix this threshold is the weak point of the approaches based on sum rules, which limits their predictivity [27].

We use the same universal effective continuum threshold in two- and three-point sum rules. This allows us to relate the value of the threshold to the pion decay constant, known experimentally. We therefore have no free numerical parameters in our analysis.

There are at least two arguments in favour of our choice of $s_0(Q^2)$:

First, we have demonstrated that it leads to the correct asymptotic behaviour of the pion form factor at $Q^2 \rightarrow \infty$. Second, we expect our approach to work better with the increase of Q^2 . We have seen that it works very well already at relatively small $Q^2 = 1\text{--}4 \text{ GeV}^2$ (recall that we have no numerical parameters to be tuned to reproduce the data). Therefore, we believe that for all $Q^2 > 1 \text{ GeV}^2$ we give reasonable predictions. However, we cannot control the accuracy of our predictions for the form factor and cannot provide any error estimates.

- We can, however, control much better the relative weights of the $O(1)$ and $O(\alpha_s)$ contributions to the form factor: their ratio is practically independent of the model for the continuum threshold and is determined to great extent by the corresponding $O(1)$ and $O(\alpha_s)$ double spectral densities. Here, our results convincingly show that the $O(\alpha_s)$ contribution to the pion form factor stays at a level below 50% at $Q^2 \leq 20 \text{ GeV}^2$ and demonstrate in a largely model-independent way that the pion form factor is mainly of nonperturbative origin up to very high Q^2 . Thus, our results definitely speak against the pQCD approach to form factors at intermediate Q^2 , referred to as Scenario A in the Introduction, and confirm Scenario B. Although obtained without any reference to the shape of the pion DA, our results indirectly restrict the pion DA at low values of the renormalization scale: For instance, convex DAs of the type of [12], close to the asymptotic one, provide the form factor compatible with the results reported here. Also a broader class of the DAs, such as, e.g., a double-humped DA with a suppressed end-point region of [22] seems to lead to the pion form factor in agreement with our results. For a conclusive clarification of this point, the analysis of both the $O(1)$ and $O(\alpha_s)$ contributions corresponding to this DA is necessary.

Finally, let us notice that the local-duality version, formulated and developed by Radyushkin and co-workers, in many cases has definite advantages compared to other versions of QCD sum rules: For instance, the standard three-point sum rules cannot go to large Q^2 because of polynomial terms, the results from light-cone sum rules depend on the light-cone distribution amplitudes. Of course, as we have already mentioned above, the numerical results for the form factor from the local-duality sum rules depend crucially on the model of the effective continuum threshold used for the calculations, but this shortcoming is shared by all versions of QCD sum rules [27]. In addition to this uncertainty, other versions of sum rules have uncertainties related to parameters not precisely known, such as the condensates and the distribution amplitudes. We therefore believe to provide the most complete analysis of the pion form factor available for the time being.

Acknowledgments

We would like to thank Alexander Bakulev and Silvano Simula for interesting discussions on the subject. D. M. was supported by the Austrian Science Fund (FWF) under projects P17692 and P20573, by RFBR project 07-02-00551, and by the Alexander von Humboldt-Stiftung. V. B. was supported by RFBR project 07-02-00417, CRDF grant Y3-P-11-05, and president grant MK-2996.2007.2.

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